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Interpretation of normal state in-plane resistivity of La-Sr-CuO superconductors**

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Abstract The temperature-dependent behaviour of the normal state in-plane resistivity of optimally doped La-Sr-CuO system is analysed within the framework of Fermi liquid description. An effective two-dimensional dynamic interaction potential is developed that incorporates the screening of holes as carriers by acoustic phonons and by plasmons. The system is treated as an ionic solid containing layers of holes as carriers and a model dielectric function is set up which fulfils the appropriate sum rules on electronic and ionic polarizabilities. The coupling strength linking holes as carriers to screened phonons is derived from the residue at the pole of interaction potential. The inelastic scattering rate due to the hole-screened phonon interaction is worked out and the Drude expression is used to estimate the temperature dependent resistivity. Zero temperature scattering rate is obtained from upper critical field and the screened plasma frequency, together with scattering rate are used to obtain zero temperature resistivity. The implications of the above analysis are discussed.

Keywords Normal state resistivity, La-Sr-CuO, screened phonons

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1. Introduction

The normal as well as superconducting state properties of high-T_c cuprates with various stacked-layer sequences offer a great challenge for theoretical explanation. As many of normal state transport properties of layered high-T_c cuprates (electrical resistivity, thermal conductivity, optical conductivity, *etc.*) are differed with those observed in typical metals, *i.e.* are unusual for a Fermi liquid. In most of the high-T_c cuprates, the inplane resistivity, ρ_{ab}, decreases linearly with decreasing temperature over a wide temperature range, while the out of plane resistivity, ρ_c, increases rapidly at low temperatures indicating the semiconducting nature [1].

We aim to study the in-plane resistivity behaviour of $La_{1.5}Sr_{0.15}CuO_4$ superconductors due to two fold reasons : A tremendous amount of work [1] has been carried out in under-doped and over-doped single crystals. It is apparent from these studies that ρ_{ab} is proportional to temperature for $(x \cong 0.15)$, whereas ρ_c is non-metallic in over- and under-doped materials [2]. A crossover [3] from 2-D transport occurs at compositions (over-doped) near the disappearance of bulk superconductivity.

Secondly, the La-based cuprates possess a relatively simple electronic structure with single conducting CuO_2 layer well separated from LaO reservoirs in comparison to that in Y-Ba-CuO systems, the presence of one-dimensional (1-D) CuO chain makes the structure little complicated. The behaviour of in-plane resistivity in over-and under-doped systems is also a significant problem, which we do not discuss in this work.

To earlier resistivity analysis of high-T_c cuprates above T_c by Lee and Read shows that inelastic scattering has contributions from electron-phonon and electron-electron scattering in a 2-D square lattice near half filling [4]. We may refer to the work of Kim *et al* who have presented a frozen phonon approach [5] for calculating the electron-phonon coupling in the presence of very strong Coulomb correlations, thereby deduce their effects on the resistivity.

The inclusion of strong Coulomb correlations is essential for creating the insulating state at half filling. An upper bound for electron-phonon coupling ($\lambda \approx 0.2$ to 0.4) is obtained from the scattering life time. The calculated mean free path from

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electron-phonon interaction shows that it is surprisingly close to the Mott-Ioffe-Regel limit [6], but it is still on the metallic site.

It is worth mentioning that Gurvitch and Fiory [7] have shown saturation effects of linear resistivity at high temperatures and estimate an upper bound of electron-phonon coupling constant $\lambda \leq 0.1$ arguing that super-conductivity in cuprates must be of non-phonon in origin. From the resistivity analysis, it is further stressed that the electron-phonon coupling strength is considerably lower despite of high-T_c values. Based on these calculations [4,5], electron-phonon contribution to ρ is nearly linear and found to account for most of the observed magnitudes. The strength of electron-phonon interaction appears to be too weak to be the primary mechanism responsible for high-T_c values and the applicability of Fermi liquid theory is questionable in layered high-T_c cuprates. The evidence [4,5,7] relating to weak electron-phonon interaction in the normal state in-plane resistivity behaviour of cuprates have provided the motivation for the present work.

The plan of the paper is as follows. In Section 2, an effective interaction potential is developed by studying the collective excitations of ions and holes in a single conducting CuO_2 layer as a limit of a periodic stack of such layers which are well separated from each other by metal-oxide layers. The random phase approximation is adopted for the polarizabilities and leads to a model dielectric function, which obeys the appropriate sum rules. The zeros of the dielectric function yield two modes, namely a low-energy acoustic plasmon and a screened phonon mode.

The electron-phonon coupling strength is calculated from the residue of the effective interaction potential $V(q, \omega)$ and later on the inelastic scattering rate. The estimated electron-phonon coupling strength favours the weak coupling strength between the hole carriers and the screened phonon mode. We estimate the electron-phonon contribution to normal state resistivity following Drude expression.

The background contribution to resistivity is calculated by using the screened plasma frequency and the zero temperature elastic scattering rates. The details of the various transport data and their results are elaborated in Section 3. The validity of Mott-Ioffe-Regel criterion for metallic conduction is also discussed. A summary and our main conclusions are presented in Section 4.

2. Essential formalism

A. *Effective interaction potential :*

The single crystal of La-based cuprate superconductors can be modelled to be an infinite array of 2D conducting CuO_2 planes which have a significant number of charge carriers and are well separated by LaO reservoirs with inter-layer distance $d_c >> d_a, d_b$. In a unit cell, the x-y plane containing a and b axes is taken to lie in the CuO_2 layer, with the c-axis lying along the z direction. To a first approximation, these layers are well separated and treated as non-interacting.

The effective interaction between each pair of holes as carriers is in our model determined by screening of the Coulomb repulsion *via* collective excitations in the layered fluid of holes (plasmons) and in the underlying ionic lattice. We begin by discussing the dielectric function. The layered electron gas systems consist of infinitely thin layers and each CuO_2 layer containing n holes as carriers are embedded in a uniform neutralising background.

The interaction energy between the hole carriers in terms of Fourier transform potential along the z-direction, namely by setting where the wave vector q_z is along perpendicular to the x-y plane as [8]

$$V_s(q, Q_z; \omega) = \frac{2\pi e^2}{q^t(q, q_z, \omega)} \sum_l \exp\left[-q_z l - q|z_l|\right] \quad (1)$$

$$\frac{2\pi e^2}{q^t(q, q_z; \omega)} S(q, q_z) \quad (2)$$

with $S(q, q_z)$ is the static form factor for the hole carriers in a single conducting CuO_2 layer sandwiched in between metal oxide layers of a unit cell and is

$$S(q, q_z) = \frac{\sinh(qd)}{\cosh(qd) - \cos(q_z d)} \quad (3)$$

The dynamic dielectric function for a single band of hole carriers is written as

$$\varepsilon(q, q_z; \omega) = \varepsilon_\infty + P(q, \omega)\, S(q, q_z)$$

with

$$P(q, \omega) = \frac{-2\pi e^2}{} \Pi(q, \omega) \quad (5)$$

in terms of polarizability function $\Pi(q, \omega)$ for a single band of hole carriers.

In view of strong anisotropy of the transport properties of cuprates [1] *i.e.*, the conduction of current carriers along and perpendicular to the CuO_2 layers are considerably different, we are mainly interested in the effective dynamic interaction between a pair of carriers in the plane at $z = 0$. Thus averaging $V_S(q, q_z; \omega)$ over q_z, we obtain the effective interaction potential which essentially explains the in-plane dynamics of the hole carriers along the conducting plane namely by setting

$$V(q, \omega) = \frac{d}{2\pi} \int_{-\pi/d}^{+\pi/d} V_S(q, q_z; \omega)\, dq_z \quad (6)$$

to get for small pd values as

$$V(q, \omega) = \frac{2\pi e^2 d}{\varepsilon_\infty} \frac{\varepsilon(q, \omega)}{\varepsilon(q, \omega)\left|\left(\varepsilon^2(q, \omega) - 1\right)^{1/2}\right.} \quad (7)$$

with $\varepsilon(q,\omega) = 1 + qdP(q,\omega)/\varepsilon_\infty$.

The high-T_c cuprates such as La-Sr-CuO with charge reservoirs and metal-oxide planes, the dopant Sr^{+2} replaces La^{+3} and effectively enters the La_2CuO_4 specimen as neutral; it then sheds its hole which are the active charge carriers. In addition, the contraction and expansion of the CuO network associated with the motions of oxygen will contribute to screening mechanism.

The effective dynamic interaction between the carriers is still given by eq. (7) : however, the dielectric function of the oxide appearing in this equation is given by the sum of contributions from the core-electrons, from the charge fluctuations in the gas of carriers and from the ionic motions. The dielectric function for a single band of hole carriers after the inclusion of ionic polarisation function becomes

$$\varepsilon(q,\omega) = 1 + qd\left[P(q,\omega) + P_i(q,\omega)\right]/\varepsilon_\infty \tag{8}$$

A rigorous numerical calculation of $P(q,\omega)$ can be made using the existing knowledge of eigenstates and eigenvalues. However, it is impossible to obtain the analytical results for the effective interaction potential that we have treated in the present investigation. The random phase approximation (RPA) forms of inter polarizability [9] have also been widely used to describe the collective excitation behaviour as well as the coupling effects in a simple manner. As we too aim to show the importance of ionic and electronic excitations we use the RPA form of $P(q,\omega)$ in eq (8)

In particular, in the long-wavelength limit $(q \to 0)$ the polarizability for a single band model of hole carriers with the 2D density n_e and the effective mass m^* yields

$$P(q,\omega) = \frac{D_1^2}{\eta^2 - \omega^2} \tag{9}$$

Here, $D_1^2 = \omega_{pe}^2 = 2\pi e^2 n_e q / m^*$ is the 2-D electron plasma frequency and $\eta^2 = q^2 v_F^2/2$, v_F being the Fermi velocity.

On the other hand, the ionic polarization function is written

$$P_i(q,\omega) = \frac{-D_2^2}{\omega^2} \tag{10}$$

where $D_2^2 = \omega_{pi}^2 = 2\pi e^2 n_i Z^2 q / M$ is the 2-D ionic plasma frequency. The sum of the ionic charge and the free electron charge is denoted by Ze, n_i is the ionic areal density and M is the mass of the CuO_2 plane.

Thus, the dynamic dielectric function becomes

$$\varepsilon(q,\omega) = 1 + \left[\alpha\Omega_1^2/\left(\eta^2 - \omega^2\right) - \alpha\Omega_2^2/\omega^2\right] \tag{11}$$

with $\alpha = qd$. The symbols Ω_1^2 and Ω_2^2 are the squares of screened electron and ion plasma frequencies.

The dynamic dielectric function in eq. (11) essentially yields the inter-layer coupled modes of the polarized waves in the layered electron gas spectrum at long wavelengths with qd as an arbitrary parameter. The evolution of main input parameters which determines the superconducting state, *i.e.*, the dynamic interaction potential $V(q,\omega)$ or equivalently the dielectric function appropriate to the specific system under consideration.

In particular, if the material is considered as a case of holes as carriers in a polar material then $\varepsilon^{-1}(q,\omega)$ is identified with the inverse longitudinal dielectric function, it is thus associated with different longitudinal modes, corresponding to exchange of phonons and plasmons, having poles at the respective modes.

The individual contributions of $\varepsilon^{-1}(q,\omega)$ from different longitudinal modes, is obtained from inversion of eq. (11) as

$$\varepsilon^{-1}(q,\omega) = 1 + \sum_{i=1}^{2} \frac{f_i(q)}{\omega^2 - \Omega_i^2}, \tag{12}$$

where Ω_i denotes the frequencies of the two longitudinal modes, to be obtained from the solution of the coupled-mode equation which follows by setting $\varepsilon(q,\omega) = 0$ at the resonance frequencies :

$$\left(\eta^2 - \Omega^2\right)\Omega^2 + \alpha\Omega_1^2\Omega^2 - \alpha\Omega_2^2\left(\eta^2 - \Omega^2\right) = 0. \tag{13}$$

Zeros of the model dielectric function which fulfil the appropriate sum rules will yield two modes of the polarized waves and the frequencies of the coupled mode in the long wavelength limit are obtained from solutions of eq. (13) as

$$2\Omega_\pm^2 = \left[\left(\Omega_1^2 + \Omega_2^2\right)\alpha + \eta^2\right] \mp \left\{\left[\left(\Omega_1^2 + \Omega_2^2\right)\alpha + \eta^2\right]^2 - 4\eta^2\Omega_2^2\alpha\right\}^{1/2}. \tag{14}$$

The above eigenfrequencies arise from mode mixing in between the acoustic phonon and plasmon modes in the layered electron gas spectrum. Further simplification is achieved by treating the second term under the square bracket as small so that, it is immediate to write

$$\Omega_+^2(q) \cong \alpha\Omega_1^2 + \eta^2 \tag{15}$$

which is the dispersion relation of an acoustic plasmon mode in harmonic approximation. The lower mode frequency is

$$\Omega_-^2 = -\frac{\alpha\Omega_2^2\eta^2}{\alpha\Omega_1^2 + \alpha\Omega_2^2 + \eta^2} \sim -\frac{\Omega_2^2\eta^2}{\Omega_1^2} \tag{16}$$

again the harmonic approximation and is a screened acoustic phonon frequency.

Thus, it is shown that the model is equivalent to a description of the system through a dielectric function in which both the

ionic and electronic contributions are summed in the random phase approximation It is convenient to rewrite the effective interaction potential in the following form [8]

$$V(q,\omega) = \frac{2\pi e^2 d}{\varepsilon_\infty} \left[1 + \frac{\Omega_+^2\left(\Omega_+^2 - \eta^2\right)}{\left(\omega^2 - \Omega_1^2\right)\left(\Omega_+^2 - \Omega^2\right)}\right.$$

$$\left. \frac{\Omega_-^2\left(\Omega^2 - \eta^2\right)}{\left(\omega^2 - \Omega_-^2\right)\left(\Omega^2 - \Omega^2\right)}\right] \quad (17)$$

As an application of the above model, the transport parameters required for the description of normal state resistivity are evaluated in the following subsection.

B. Normal state resistivity :

The normal state resistivity is well explained from the linearized Bloch Boltzmann transport equation where the electron-phonon interaction is dominant [10]. We first deduce the coupling strength $\gamma_-(q)$ linking hole carriers to the $\Omega_-(q)$ mode from the residue of $V(q,\omega)$ as

$$\frac{\gamma_-|^2}{\Omega} = \frac{2\pi e^2 d}{\varepsilon_\infty} \frac{\alpha\left(\Omega_1^2 + \Omega_2^2\right) - \eta^2}{\alpha\left(2\Omega_1^2 - \Omega_2^2\right) - \eta^2}. \quad (18)$$

The temperature dependent part of the normal state resistivity is included in the inelastic scattering rate due to electron- phonon interaction The mean time to absorb or emit a phonon of energy $\hbar\Omega_-(q)$ is expressed as [9]

$$\tau_{e-ph}^{-1} = \frac{2\pi}{\hbar}\int_0^{2k_f} \frac{d^2q}{(2\pi)^2} |\gamma_-|^2 (1 - \cos\theta)$$

$$\times\left[n_q \delta\left(\varepsilon_{k+q} - \varepsilon_k - \hbar\Omega_-\right) + \left(n_q + 1\right)\delta\left(\varepsilon_{k+q} - \varepsilon_k - \hbar\Omega_-\right)\right] \quad (19)$$

Here, n_q is the phonon occupancy factor and for high temperatures, n_q is reduced to $k_B T/\hbar\Omega_-$. Polar angle θ is in between k and $k+q$ Integrating over the angle ϕ between q and k, we obtain

$$\frac{1}{\tau_{e-ph}} = \frac{m^* k_B T}{\pi\hbar^3 k_f^3}\int_0^{2k_f} \frac{|\gamma_-|}{\Omega_-^2} \frac{q^2 dq}{\left[1 - (q/2k_f)^2\right]^{1/2}}. \quad (20)$$

Integrating over the two-dimensional wavevector (q) with the use of eq. (20), we obtain the inelastic scattering rate as

$$\frac{1}{\tau_{e-ph}} = \frac{2\pi e^2 m^* k_B T}{\hbar^3 k_F \varepsilon_\infty} \frac{\left(\Omega_1^2 - \Omega_2^2\right)d/2k_F - v_F^2}{\left(\Omega_1^2 - 2\Omega_2^2\right)d/2k_F - v_F^2}. \quad (21)$$

Usually, δ function fixes the angle between the energies in 2-D case and hence linear temperature dependence is obtained of the inelastic scattering rate. The normal state resistivity due to electron-phonon interaction in terms of scattering time is expressed as

$$\rho_{e-ph} = \left[m^*/n_c e^2\right]\left[1/\tau_{e-ph}\right]. \quad (22)$$

The temperature dependent part of the resistivity is obtained through electron-phonon scattering within the framework of layered electron gas spectrum. In addition to this scattering mechanism, other scattering mechanisms are also possible. The electrons also scatter off impurities, defects and disordered regions and thereby giving rise to a temperature dependent contribution.

Knowledge of zero temperature scattering rate and the plasma frequency will allow us to have an independent estimation of zero temperature resistivity [11]. In a true sense, the Fermi liquid based schemes are meaningful at absolute zero and for the reason the T_c is rather low, we may use the term loosely to characterise the superconducting state where the resistivity is minimum on the scale of measurements.

The in-plane resistivity is now modelled as the sum of zero temperature resistivity and the temperature dependent part, which is indebted to the electron-phonon scattering process Hence,

$$\rho_{total}(T) = \rho(0) + \rho_{e-ph}(T). \quad (23)$$

Using the development expressions, for the normal state inplane resistivity, the numerical results on various model parameters are presented along with discussions in the following section.

3. Results and discussion

In the present work, we have emphasised on the screened phonon contribution for the description of normal state in-plane resistivity of optimizes doped La-Sr-CuO superconductors. For this purpose, realistic values of some physical parameters are derived from experimental data as follows. For a layered stacking sequence well separated by an average spacing, $d(= c/2)$, the effective mass of the carriers along the conducting CuO_2 plane is obtained from the electronic specific heat coefficient γ using the relation, $m^* = 3\hbar^2\gamma d/\pi k_B^2$. Taking interplanar distance, d = 6.6Å and γ = 5 mJ/mol/K^2 from the specific heat measurement, $m^* \equiv 4m_e$ is obtained.

The behaviour of the multilayer systems critically depends only on the planar carrier density and the spacings between the layers, and not on the form of single particle distribution function, the layer thickness, or the number of layers. For a stack of 2-D conducting planes well-separated by an average spacing d, the condition for optimised pairing [11] infers the 2-D charge carrier density and follows $n_c d^2 = 1$ to obtain $n_c \equiv 2.3\times10^{14}$ cm^{-2}

The other parameters of the electrons are Fermi velocity $v_F \cong 1.1 \times 10^7$ cm s^{-1} and the screened plasma frequency is about 0.81 eV.

Turning to the parameters related to the ions, a reasonable value of the background dielectric constant is $\varepsilon_\infty \cong 4.5$. The ionic areal density is the reciprocal of the unit cell projection perpendicular to the c-axis and is $n_i \cong 7 \times 10^{14}$ cm^{-2} from the relation $n_i = 1/ab$. The lattice parameters [12] $a = b = 3.7793$Å and $c = 13.226$ Å (for $x = 0.15$) are used for the above purpose. The charge distribution in CuO_2 plane includes the ionic charge and free electron charge, the ionic charge Ze being –2e.

The mass of a unit cell is $M = 12.77$ amu. The attractive range of the interaction potential are important for small qd values and the layering effects are insignificant for higher qd values as the mode merge from each other. Since our focus is on the long-wavelength ($q \to 0$) behaviour of the collective modes, we have considered the long-wavelength limit in the intermediate ($\alpha = qd = 0.1$) coupling limit.

For the sake of simplicity, one oscillator for holes as carriers and ions are considered. Since the system contains 7 atoms in a unit cell, various modes developed can be understood by adding more oscillators in the model dielectric function. With the above values for the input parameters, the upper and lower mode frequencies $\hbar\Omega_+$, and $\hbar\Omega_-$ are then estimated. The inter-layer acoustic plasmon is estimated as 0.46 eV and the frequency of the dressed phonons with electronic excitations is about 17.6 meV. With the use of lattice parameter a, and the screened phonon energy $\hbar\Omega_-$, the long-wavelength sound velocity $v_s (= ak_B\Omega_- / \hbar\pi)$ is estimated as 4×10^5 cm sec^{-1}.

The scattering of charge carriers at Fermi surface is considered for all possible values of scattering angle θ. The 2-D wave vector $q(= 2k_F \sin\theta)$ can, therefore, take maximum values up to $2k_F$. The scattering processes involve $2k_F$ scattering across the Fermi surface and the scattering rate due to electron-phonon interaction is directly proportional to the temperature. We estimate τ_{e-ph} as 2.4×10^{-14} sec at $T = 100$ K for the present La-based system. The reflectance spectra data [13] on single crystal thin films of La-Sr-CuO for the best fits to the Drude model yields the scattering rate as 0.86×10^{-14} sec. It is further stressed that τ increases from underdoped to optimally doped and in the over-doped region ($\lambda \approx 0.3$, $\tau = 1.3 \times 10^{14}$ sec) and it becomes almost Sr concentration independent.

We may also refer to the work of Kim *et al* [5] who have evaluated the electron-phonon contribution to the transport life time by solving the Boltzmann transport equation variationally and the Eliashberg function $\alpha^2 F(\omega)$. The inelastic scattering rate from transport data is approximately 5×10^{-14} sec at $T = 100$K in $La_{1.85}Sr_{0.15}CuO_4$ system. The present model demonstrates that the scattering life time from the electron-phonon interaction is roughly 3 times larger than the reflectance data [13] and is smaller by a factor of 2 from the theoretical estimation [5]. It is expected that due to mass enhancement ($m^* \cong 4m_e$) of carriers in La-Sr-Cuo, the inelastic scattering rate exceeds k_BT and we estimate $r^{-1} \cong 3k_BT$. We recognize the low frequency screened phonons as the primary source of transport lifetime.

In typical metals, the inelastic scattering rate is related to the electron-phonon coupling strength through $r^{-1} \approx 2\pi\lambda_{tr}k_BT$. For the estimated τ^{-1}_{e-ph} of about 0.42×10^{14} sec^{-1} at $T = 100$K, we find $\lambda_{tr} = 0.5$. The striking feature observed in the resistivity saturation [10] at high temperatures allows an estimation of the upper bound on λ as 0.1. The authors [7] pointed out that the weakly coupled phonon could not yield high-T_c and superconductivity must be non-phonon in nature for cuprates. We shall also refer to the work of Kim *et al* [5] who have proposed that upper bound of inelastic electron-phonon coupling constant is roughly 0.2 ~ 0.4 in La-Sr-CuO systems which is significantly different from the previous estimates of ranges [14] from 0.65 to 2.0. These numbers are in excess of the upper bound imposed by *ac* conductivity Drude fits. Precisely, it is commented that the electron-phonon coupling appears to be too weak for the primary mechanism although their participation cannot be ruled out and we therefore address this issue

The electron-phonon coupling strength determines well the superconducting T_c in Eliashberg theory. Using the deduced λ_{tr} value as 0.5, one could not meet T_c of 40K and higher. There are two points to be made here. First, the electron-phonon coupling constant (λ) appears low and hence reflects the weak interaction between hole carriers and the screened phonons. Furthermore, the inelastic scattering rate is several time k_BT and is consistent with the inelastic neutron scattering data, as the gap ratio $2\Delta(0) \approx 6k_BT_c$ and $\Delta(0) \approx \tau^{-1}_{in}(T_c)$. It seems that the deduced coupling constant does not account for high-T_c value with phonon mechanism alone and essentiality of incorporation of collective excitations in order to predict the superconducting T_c. Secondly, for materials with nested Fermi surface, the phonon scattering involves $2k_F$ scattering across the Fermi surface, which dissipates momentum, and hence there exists a significant difference in λ_{tr} and λ_{in}. It is thus believed that for any operative mechanism in copper oxides, the normal state transport data provides a constraint and must be recognised in first step in revealing the superconducting state parameters.

A temperature dependence of in-plane resistivity of $La_{1.85}Sr_{0.15}CuO_4$ in Figure 1 is explained by the zero temperature limited resistivity and contribution from inelastic hole carrier-scattering. The calculated values of in-plane resistivity are although smaller than the single crystal data [3] but lies in the same range qualitatively. We believe that the upper critical magnetic field, $H_{c2}(0)$, and the screened plasma frequency represents a good set of the physical parameters for the estimation of zero temperature resistivity as $\rho(0) = 0.058$ mΩ cm. This is consistent with single crystal data [3].

We find that calculated $\rho[=\rho(0)+\rho_{e-ph}(T)]$ does not exceed the measured ρ_{ab} in whole temperature range ($T_c < T <$ 300). It is natural to comment that the displayed temperature dependence of normal state resistivity (ρ) is well explained qualitatively by inter-layer screened phonon mechanism within the fremework of layered electron gas approach

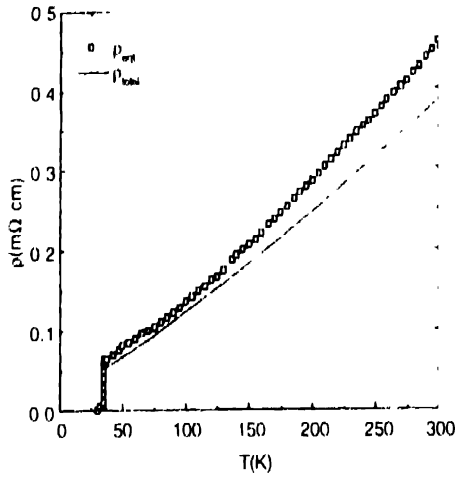

Figure 1. Variation of normal state in-plane resistivity with $T(K)$ for $La_{1.85}Sr_{0.15}CuO_4$ along with the single crystal [3] data(0)

4. Conclusion

In this communication, we have devoted our efforts to study the normal state in-plane resistivity of $La_{1.85}Sr_{0.15}CuO_4$ superconductors within the framework of Fermi liquid description. The model dynamic dielectric function associated with polarizability of hole as carriers and ions in the conducting plane properly incorporate the form factor of the stacked layer sequence. Confining the conduction along the plane with suitable transformation, an effective two dimensional interaction potential is set up for the carriers in the CuO_2 layers. The hole-screened phonon (low frequency) coupling constant is worked out from the residues at the poles of the interaction potential and hence the inelastic scattering rate

The appropriateness of the developed approach is mainly due to the attention on CuO_2 plane in the layered fluid of holes and in the underlying ionic lattice. We believed that the Fermi curve has nested states and as a result, we dealt with $2k_F$ scattering across the Fermi surface. The inelastic scattering rate due to electron-phonon interaction is then estimated roughly $3k_BT$ due to the mass enhancement of hole as carriers. We estimate the coupling constant as one half, which points towards the weak inter-layer, hole carrier-phonon interactions. The backward scattering of holes in the nested Fermi surface enhances the transport coupling constant by a factor of two and the significant difference in λ_{in} and λ_{tr} puts constraints on any operative superconducting mechanism.

An important feature of the above analysis is the validity of the Mott-Ioffe-Regel criterion for metallic conduction in cuprates The mean free path at $T = 100$ K is several times larger than the copper oxide bond-length. A comparable value of mean free path with zero temperature coherence length essentially points towards the clean limit of optimally doped La-based cuprates. The quantitative estimation of zero temperature scattering rates from upper critical magnetic field indicates smaller plasma frequency and the Fermi velocity. Further, we find that the magnitude of inelastic scattering rate increases with temperature The zero temperature resistivity thus estimated with the above said informations is consistent with single crystal data.

In conclusion, the large magnitude of resistivity at room temperature as well as the temperature dependence of optimised La-based cuprates is well understood with low frequency screened phonon mechanism as arises from inter-layer interactions in the framework of Fermi liquid description. The present model can easily be extended to the other members of the family of cuprate superconductors having a larger number of CuO_2 layers. Multiple layers can be incorporated in the effective interaction potential and one may expect additional excitation branches.

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